Preprint

Lawrence Berkeley Laboratory
UNIVERSITY OF CALIFORNIA

Physics, Computer Science & Mathematics Division

Submitted to Physics Letters B

SOFT-GLUON EFFECTS IN NONLEPTONIC DECAYS OF CHARMED MESONS

Ken-ichi Shizuya

November 1980

Prepared for the U.S. Department of Energy under Contract W-7405-ENG-48

DISCLAIMER

SOFT-GLUON EFFECTS IN NONLEPTONIC DECAYS OF CHARMED MESONS

Ken-ichi Shizuya

Lawrence Berkeley Laboratory
University of California
Berkeley, California 94720

ABSTRACT

Soft-gluon effects in nonleptonic decays of D and F mesons are studied nonperturbatively by use of a QCD multipole expansion. For reasonable values of D-meson bound-state parameters, the soft-gluon effects lead to a significant difference in the lifetimes of the D^0 and D^+ mesons.

Recent experiments [1] have reported an appreciable difference in the lifetimes of the D^0 and D^+ mesons, with $\tau(D^+)/\tau(D^0) = 3 \sim 10$. A simple picture of charmed meson decays, based on the charm-quark decay process $c \to su\bar{d}$ or $s\nu\bar{\ell}$ (fig. 1(a)), predicts equal lifetimes for D^0, D^+ and F^+ mesons. The observed lifetime difference suggests significant enhancements of nonleptonic D^0 decays by some dynamical mechanisms [2].

The W-exchange process ("quark-annihilation" process), as depicted in fig. 1(b), contributes solely to D^0 decays. A helicity factor $(m_s/M_D)^2$ and the small probability for quark-pair annihilation in the D meson combine to make its contribution negligibly small. A number of authors [3-5], however, have pointed out that the presence of gluons may be crucial to the removal of helicity suppression which otherwise is inherent in quark-annihilation processes. Single-hard-gluon emission from the D^0 meson [3-5], as depicted in fig.2(a), serves to remove helicity suppression of the accompanying weak decay and enhances the D^0 decay rate; however, such short-distance QCD effects alone seem to be too small to account for the D^0-D^+ lifetime difference. Soft gluons inside (or surrounding) the D meson are equally likely sources of enhancement for nonleptonic D^0 decays. Because of the nonperturbative nature of soft-gluon interactions, however, estimates of their effects have so far been purely phenomenological [4].

The purpose of this paper is to present a dynamical calculation of the soft-gluon effects on D^0-decay enhancements. We study the soft-gluon effects nonperturbatively by use of a QCD multipole expansion [6–8] and relate them to the vacuum expectation value

$$\mathscr{V} = \langle 0 |(\alpha_S/\pi) F_{\mu\nu}[A]^2 | 0 \rangle \sim 0.012\ \mathrm{GeV}^4, \qquad (1)$$

whose magnitude is phenomenologically known from the charmonium sum rules of Shifman, Vainshtein and Zakharov [9]. This matrix element provides a measure of the nonperturbative soft-gluon fluctuations residing in QCD color-confinement mechanisms.

Figure 2(b) represents the process we consider. A hard gluon in fig. 2(a) is replaced by a collection of soft gluons here. (What is meant by a collection of soft gluons will be made clearer below.) We treat the D^0 meson as a nonrelativistic bound state[F1] of c and $\bar{u}$ constituent quarks of mass $m_c \approx 1.65$ GeV and $m_u \approx 0.34$ GeV. Soft-gluon emission from the final quarks is not taken into account since it is not directly related to the removal of helicity suppression factors. Both the motion and the soft-gluon interactions of the c quark are ignored in what follows; they are suppressed by a ratio m_u/m_c in the decay amplitude as compared with those of the $\bar{u}$ quark.

Let us expand the soft-gluon field around the c quark into multipoles. The multipole terms are cast into a manifestly gauge-invariant form by a suitable transformation [7,8]. The interaction

of the u quark with the gluon field is described by the Hamiltonian

$$\mathscr{H}^{QCD} = g\bar{u}(x)\left[\gamma^0 \vec{x}\cdot\vec{E}^a(\vec{0}) + \frac{1}{2m_u}(\vec{\sigma} + \vec{L})\cdot\vec{H}^a(\vec{0})\right]\left(\frac{1}{2}\lambda^a\right)u(x) + \ldots, \qquad (2)$$

where $u(x)$ is the u quark field at position $\vec{x}$, $E^{ka}(\vec{0}) = \left(F^{k0}[A(\vec{0})]\right)^a$ and $H^{ka}(\vec{0}) = -\frac{1}{2}\epsilon^{k\ell m}F^a_{\ell m}[A(\vec{0})]$ ($F_{\mu\nu}[A] = \partial_\mu A_\nu - \partial_\nu A_\mu + gA_\mu \times A_\nu$) are the soft-gluon fields defined at the c quark position ($\vec{x} = \vec{0}$), and $\vec{L} = \vec{x}\times\vec{p}$ ($p^k = -i\partial/\partial x^k$) is the angular momentum of the u-quark motion. Only the color-E1 and the color-M1 interactions are shown. The color-Coulomb interaction does not act on a color-singlet $c\bar{u}$ system. Higher multipole interactions are not included. Note that the $\vec{L}$ term in (2) does not contribute to an S-wave state.

The initial D^0 meson is a color-singlet (1), ${}^1S_0$ ($c\bar{u}$) bound state. The color-E1 interaction turns the D^0 meson into a color-octet (8), P-wave $c\bar{u}$ state while the color-M1 interaction changes both the color and spin of the $c\bar{u}$ system; namely,

$$(1, {}^1S_0) \begin{array}{l} \nearrow (8, {}^1P_1) + \mathscr{G}\ (\mathrm{E1}), \\ \searrow (8, {}^3S_1) + \mathscr{G}\ (\mathrm{M1}). \end{array} \qquad (3)$$

Here $\mathscr{G}$ stands collectively for the color-octet states that are reached by either a color-E1 or a color-M1 soft-gluon interaction; we call $\mathscr{G}$ the "soft gluons".

The virtual color-octet $c\bar{u}$ systems have $J^P = 1^+$ or $J^P = 1^-$ so that their subsequent weak decays are free of helicity suppression. The $\Delta S = \Delta C = -1$ nonleptonic weak Hamiltonian is given by (we set the Cabbibo angle $\theta_C = 0$)

$$\mathcal{H}^W = \frac{4G}{\sqrt{2}}\left[\left(\frac{1}{N} f_1 + f_2\right) (\bar{s}d)_L(\bar{u}c)_L + \frac{1}{2} f_1(\bar{s}\lambda^a d)_L(\bar{u}\lambda^a c)_L\right] + \text{h.c.}, \tag{4}$$

where Lorentz as well as color indices have been suppressed in usual fashion; $(\bar{u}c)_L \equiv \bar{u}\gamma^\mu\frac{1}{2}(1-\gamma_5)c$, etc. This is a Fierz-reordered form (suitable for D^0 decays) of the conventional weak Hamiltonian. In the above, $N = 3$ for color SU(3); we use SU(N) notation to keep track of color factors. Hard-gluon exchange corrections to the weak Hamiltonian [10] change the coefficients f_1 and f_2 from their zeroth-order values ($f_1 = 1$ and $f_2 = 0$) to $f_1 \sim 1.42$ and $f_2 \sim -0.74$.

Let us denote the D^0-meson wave function in momentum space by $\psi(\vec{p})$ normalized so that $(2\pi)^{-3}\int d^3p|\psi(\vec{p})|^2 = 1$, where $\vec{p}$ is the momentum of the $\bar{u}$ quark in the D^0 meson. Note that the S-wave function $\psi(\vec{p})$ is a function of $|\vec{p}|$. As usual, quark spinors are approximated by free quark spinors. The virtual $c\bar{u}$ states can be decomposed into spin-zero and spin-one components by use of appropriate spin-projection operators. The old-fashioned perturbation theory leads to the following amplitude for the process in fig. 2(b).

$$\mathcal{F} = \sqrt{\frac{M_D}{N}}\int\frac{d^3p}{(2\pi)^3}\frac{1}{M_D-\varepsilon} L^a_\mu(\vec{q}_1,\vec{q}_2)J^{\mu a}(\vec{p}, \vec{\omega})\psi(\vec{p}), \tag{5}$$

where ε is the energy of the intermediate $(c\bar{u})_8 + \mathcal{G}$ state, $\vec{\omega}$ is the momentum of the "soft-gluon" state $|\mathcal{G}>$ and $L^a_\mu(\vec{q}_1,\vec{q}_2)$ refers to the weak current of the light quarks

$$L^a_\mu(\vec{q}_1,\vec{q}_2) = \sqrt{2}G(\tfrac{1}{2}f_1)\bar{u}_s(q_1)\gamma_\mu\lambda^a(1-\gamma_5)v_d(q_2). \tag{6}$$

For the color-E1 interaction $J^{\mu a}(\vec{p},\vec{\omega})$ is given by

$$J^{\mu a}(E1) = g <\mathcal{G}|E^{ka}(\vec{0})|0> \begin{bmatrix} x^k \\ p^\ell x^k/(2m_u) \end{bmatrix} \begin{matrix} (\mu=0) \\ (\mu=\ell) \end{matrix} \tag{7}$$

while for the color-M1 interaction

$$J^{\mu a}(M1) = -\frac{1}{2}\frac{g}{m_u} <\mathcal{G}|H^{ka}(\vec{0})|0> \begin{bmatrix} p^k/(2m_u) \\ \delta^{k\ell} + i\varepsilon^{k\ell j}p^j/(2m_u) \end{bmatrix} \begin{matrix} (\mu=0) \\ (\mu=\ell) \end{matrix}, \tag{8}$$

where terms quadratic in $\vec{p}$ or higher as well as those[F2] that vanish as $\vec{\omega} \to 0$ have been omitted.

Soft gluons are strongly interacting and need to be treated nonperturbatively. The decay rate, e.g. for the process with the color-M1 interaction, involves a sum over the soft-gluon states $|\mathcal{G}>$ of the form

$$\sum_{\mathcal{G}} <0|H^{ib}(\vec{0})|\mathcal{G}> <\mathcal{G}|H^{ja}(\vec{0})|0> \frac{1}{(M_D-\varepsilon)^2}(\cdots) \quad . \tag{9}$$

The energy denominator $M_D - \varepsilon$ represents the energy difference between the initial D^0-meson state and the virtual $(c\bar{u})_8 + \mathcal{G}$ state. It will be reasonable to suppose that, owing to color confinement, the virtual color-octet $c\bar{u}$ system has higher energy than the initial D^0-meson state and that this energy difference does not vanish even for very soft gluons ($\vec{\omega} \to 0$). [For example, this is the case for mesons bound by a one-gluon-exchange potential, which is repulsive between a color-octet quark-antiquark pair.] This energy difference will be of the order of $100 \sim 200$ MeV, a typical scale related to confinement; this point will be discussed later. With these in mind, we approximate the soft-gluon sum in (9) as follows

$$\text{(i)}\quad \sum_{\mathcal{G}} |\mathcal{G}> (\varepsilon - M_D)^{-1} <\mathcal{G}| \sim (\Delta\varepsilon)^{-1} 1. \qquad (10)$$

Namely, the energy denominator $\varepsilon - M_D$ is replaced by a constant value $\Delta\varepsilon$ typical for the soft-gluon reaction we consider. It is understood that eq. (10) is inserted between soft-gluon states (i.e. excluding hard-gluon states); accordingly we set $\vec{\omega} \sim 0$ in $(\cdots)$ of eq. (9).

(ii) The soft-gluon world would exhibit approximate Lorentz invariance relative to the wavelengths of color fluctuations in it. By making use of this invariance, eq. (9), which is now rewritten as

$$< 0|H^{ib}(\vec{0})H^{ja}(\vec{0})|0 > (\Delta\varepsilon)^{-2}(\cdots), \qquad (11)$$

is related to $< 0|F^2_{\mu\nu}|0>$ in eq. (1).

An argument in favor of these approximations[F3] is made as follows: The phenomenological value for $\mathcal{V} = < 0|(\alpha_S/\pi)F^2_{\mu\nu}|0 >$ in eq. (1) represents nonperturbative (long-distance) QCD effects[F4]; namely, this matrix element is predominantly saturated by the long-distance color fluctuations of confinement physics. Each of the above approximations relies upon and could be justified by this dominance of long-distance dynamics in the phenomenological value of $\mathcal{V}$.

The value of the unknown quantity $\Delta\varepsilon$ may be estimated in the following way. The annihilation of the virtual color-octet $c\bar{u}$ system by the weak current occurs in a small domain of the size $1/m_c$, the Compton wavelength of the c quark. Correspondingly, $\Delta\varepsilon$ will be relatively sensitive to the short-distance structure (e.g. the spin-dependent part) of the $c\bar{u}$ binding potential. Although the virtual $c\bar{u}$ system is in color-octet, the $(c\bar{u})_8 + \mathcal{G}$ state as a whole is a color-singlet and presumably is still in a bound state [of spatial spread of the order of $1/\Delta\varepsilon$]. Namely, the virtual state may be picturized as a spin-one D-meson system with a (color-octet) gluon cloud; and the total energy of this meson system may not be drastically affected by the spatially small color concentration of the $c\bar{u}$ system. If this picture is adequate, the $^3S_1 - {}^1S_0$ fine structure of the D-meson system provides a reasonable guess for $\Delta\varepsilon$ (for the color-M1 process):

$$\Delta\varepsilon \sim M(D^*) - M(D) \approx 140 \text{ MeV}. \qquad (12)$$

A naive estimate of the "binding energy" of the D^0 meson gives another measure of $\Delta\varepsilon$:

$$\Delta\varepsilon \sim m_c + m_u - M_D \approx 120\ \text{MeV}\ . \tag{13}$$

With the abovementioned approximations, the amplitude (9) simplifies. In particular, since $x^j = i\partial/\partial p^j$ in $\vec{p}$ space,

$$\int d^3p\ p^\ell x^j \psi(\vec{p}) = -\ i\delta^{\ell j}\int d^3p\psi(\vec{p}) \tag{14}$$

by an integration by parts.[F5] In the nonrelativistic approximation, the wave function at the origin $\vec{x} = 0$, $\phi(\vec{0}) = (2\pi)^{-3}\int d^3p\psi(\vec{p})$, is related to the D-meson decay constant f_D so that

$$f_D = 2(N/M_D)^{1/2}|\phi(\vec{0})|.$$

Apart from an unobservable overall phase factor, the amplitude $\mathcal{F}$ is now rewritten as

$$\mathcal{F} \simeq (4m_u\Delta\varepsilon N)^{-1} g f_D M_D L^a_\ell(\vec{q}_1,\ \vec{q}_2) \langle \mathcal{G}| H^{\ell a}(\vec{0}) + iE^{\ell a}(\vec{0})|0\rangle, \tag{15}$$

where we have taken the same $\Delta\varepsilon$ for the color-M1 and color-E1 transitions (although $\Delta\varepsilon$ may well be different in the two cases). In the calculation of the decay rate, we use the relation [9]

$$\langle 0|H^a_k H^a_\ell|0\rangle = -\langle 0|E^a_k E^a_\ell|0\ \rangle = \frac{1}{12}\delta^{k\ell}\langle 0|F^2_{\mu\nu}|0\rangle, \tag{16}$$

in accordance with the approximations explained earlier. Note that E^a_k is antihermitian in the nonperturbative QCD vacuum. The contributions of the color-E1 and the color-M1 processes to the decay rate turn out to be of equal magnitude. The decay rate $\Gamma^{sg} = \Gamma^{sg}(E1 + M1)$ for the soft-gluon process is given by

$$\Gamma^{sg} = \frac{G^2 f_D^2 M_D^3}{8\pi}\ \frac{8}{3}\left(\frac{f_1}{N}\right)^2\left(\frac{\pi}{4m_u\Delta\varepsilon}\right)^2 \mathcal{V}\ , \tag{17}$$

where $\mathcal{V} = \langle 0|\frac{\alpha_S}{\pi}\ F^2_{\mu\nu}|0\rangle$ (eq. (1)), and the light quark masses in the final state have been neglected. On the other hand, the c-quark decay process in fig. 1(a) leads to equal D^0 and D^+ decay rates

$$\Gamma^c = (2 + Nh)G^2 m_c^5/(192\pi^3)\ , \tag{18}$$

where the factor 2 is for leptons and the factor $h \equiv f_1^2 + f_2^2 + (2/N)f_1 f_2 \sim 1.8$ involves the hard-gluon exchange effect. As before, the final quark masses have been neglected. The relative importance of the soft-gluon process to the quark-decay process is seen from the ratio

$$R = \Gamma^{sg}(E1 + M1)/\Gamma^c$$

$$= \frac{f_1^2}{(2+Nh)}\ \left(\frac{2\pi^2}{N}\right)^2\left(\frac{M_D}{m_c}\right)^3\left(\frac{f_D}{m_c}\right)^2\ \frac{\mathcal{V}}{(m_u\Delta\varepsilon)^2}\ . \tag{19}$$

Substituting the numerical values quoted earlier for m_u, m_c and $\mathcal{V}$ yields

$$R \approx 0.7 \times (f_D/\Delta\varepsilon)^2\ . \tag{20}$$

A reasonable estimate for f_D, based on an empirical scaling law for the observed lepton-pair decay rates of vector mesons, gives [5]

$$f_D/\sqrt{2} \sim 150 \text{ MeV} . \qquad (21)$$

With this value for f_D and $\Delta\varepsilon \sim 140$ MeV (120 MeV), the present soft-gluon process leads to the lifetime difference

$$\tau(D^+)/\tau(D^0) = 1 + R \sim 2.5 \ (3.0) . \qquad (22)$$

This result indicates that the soft-gluon effect by itself could account for a significant portion of the lifetime difference between the D^+ and D^0 mesons.

The decay rate Γ^{hg} for the single-hard-gluon emission process in fig. 2(a), evaluated perturbatively [3,5], is smaller than the soft-gluon effect Γ^{sg}:

$$\Gamma^{hg}/\Gamma^{sg} = \frac{2}{3}\,(\alpha_S/\pi^3)\ (M_D\Delta\varepsilon)^2/\mathscr{V}$$

$$\approx 0.02\,\alpha_S \times (\Delta\varepsilon/60 \text{ MeV})^2 , \qquad (23)$$

where the α_S is the coupling constant characterizing the hard-gluon emission process. In D^0 decays, therefore, the hard-gluon emission effect is one order of magnitude smaller than the soft-gluon effect.

Quantitatively the present analysis is a crude estimate because of uncertainties involved in $\Delta\varepsilon$ and f_D and of the approximation scheme employed. It is conceivable that higher-multipole terms, especially those of soft gluons coupled to gluon exchanges between the constituent quarks in the D meson [8], are non negligible for the D-meson system (although they are less important for heavy quarkonium systems). Qualitatively, however, the abovementioned conclusions of the present analysis will, we expect, survive a more detailed analysis.

Soft-gluon effects are expected to be important in other heavy-meson decays as well.

(1) Charmed F-meson decays. The present analysis developed for the D^0 meson is carried over to the F^+ meson by interchanging u quarks $\leftrightarrow$ s quarks and $f_1 \leftrightarrow f_2$. It is owing to the hard-gluon exchange corrections to the weak Hamiltonian that color-octet $c\bar{s}$ systems are annihilated into u and $\bar{d}$ quarks via the process in Fig. 2(b); with either a single color-E1 or color-M1 soft-gluon interaction, their annihilation into $\nu\bar{\ell}$ is still forbidden by color mismatch. The ratio of the soft-gluon corrections Γ^{sg} for the F^+ and D^0 mesons is given by

$$\frac{\Gamma^{sg}(F^+)}{\Gamma^{sg}(D^0)} = \left(\frac{f_2}{f_1}\right)^2 \left(\frac{f_F}{f_D}\right)^2 \left(\frac{m_u\Delta\varepsilon_D}{m_s\Delta\varepsilon_F}\right)^2 \left(\frac{M_F}{M_D}\right)^3 . \qquad (24)$$

With $M_F \approx 2.03$ GeV, $m_s \approx 0.54$ GeV (constituent quark mass), $f_F/f_D \sim 1.4$ (an estimate based on the scaling law) and

$\Delta\epsilon_F \sim M(F^*) - M(F) \sim 110$ MeV, one obtains the estimate

$$\Gamma^{sg}(F^+)/\Gamma^{sg}(D^0) \sim 0.4 \ . \qquad (25)$$

This indicates that the F-meson lifetime is somewhere between those of the D^0 and D^+ mesons:

$$\tau(D^0) < \tau(F^+) < \tau(D^+) \qquad (26)$$

(2) For heavy mesons containing b or t quarks, the difference in the lifetimes of the neutral and charged members will not be as prominent as in the case of D decays. The enhancement of annihilation processes by the soft-gluon effect diminishes rapidly with increasing heavy-quark mass: The ratio R in eq. (19) decreases like $1/m_c^3$ as $m_c \to \infty$ (since $f_D \propto m_c^{-1/2}$). Correspondingly, in B-meson decays this ratio will presumably be more than one order-of-magnitude smaller than in D decays. As seen from eq. (23), the soft-gluon effect and the hard-gluon emission effect are comparable in B-meson decays. For sufficiently heavy mesons, nonleptonic enhancements will be dominated by short-distance mechanisms, hard-gluon emission from the initial and final quarks.

ACKNOWLEDGMENTS

I would like to thank M. Chanowitz and M. Suzuki for helpful comments and a careful reading of the manuscript and S. Nussinov for useful discussions. This work was supported by the High Energy Physics Division of the U. S. Department of Energy under contract no. W-7405-ENG-48.

FIGURE CAPTIONS

Figure 1. D^0 decays. (a) charm-quark decay process. The $\bar{u}$ quark acts as a spectator. (b) quark-annihilation process. shaded blobs represent W-boson exchanges with hard-gluon exchange corrections.

Figure 2. Quark-annihilation process with emission of gluons from the D^0 meson. (a) hard-gluon emission. (b) soft-gluon emission. The dashed line refers to a virtual state.

FOOTNOTES

F1 The nonrelativistic description of the D-meson system, though not as reliable as that of $c\bar{c}$ charmonium systems, has certain phenomenological success; see, e.g., ref.[10]. The nonrelativistic picture improves somewhat for $c\bar{s}$ systems.

F2 Terms that depend on $\vec{\omega}$ correspond to higher-multipole terms (other than color-E1 and color-M1 terms).

F3 Similar approximations have previously been used for heavy-quarkonium systems by Voloshin [6].

F4 In standard perturbation theory, the contribution of very soft gluons to $\langle 0|F_{\mu\nu}^2|0\rangle$ is vanishingly small at least to lowest order because of derivative coupling. Hard gluons mainly contribute to the renormalization of the operator $\alpha_S F_{\mu\nu}^2$ (which is a renormalization-group invariant in the absence of quarks). The nonvanishing value of $\langle 0|\alpha_S F_{\mu\nu}^2|0\rangle$, being a long-wavelength phenomenon, is considered to be a consequence of strong soft-gluon interactions.

F5 The surface term in eq. (14) vanishes for a large class of quark-binding potentials, including the Coulomb and the harmonic oscillator potentials.

REFERENCES

[1] J. Kirby, Proc. Intern. Symp. on Lepton and Photon Interaction at High Energies (Fermilab, August 1979); V. Lüth, ibid; J. Prentice, ibid; W. Bacino et al., Phys. Rev. Letters 45 (1980) 329.

[2] For a recent review on charmed-meson decays, see M. Chanowitz, High Energy e^+e^- Interactions (Vanderbilt, 1980), AIP Conf. Proc. No. 62 (American Inst. of Physics, NY, 1980) p. 48.

[3] M. Bander, D. Silverman and A. Soni, Phys. Rev. Letters 44 (1980) 7.

[4] H. Fritzsch and P. Minkowski, Phys. Letters 90B (1980) 455.

[5] M. Suzuki, UC-Berkeley preprint UCB-PTH-8014, to appear in Nucl. Phys. B.

[6] K. Gottfried, Phys. Rev. Letters 40 (1978) 598; M. B. Voloshin, Nucl. Phys. B 154 (1979) 365; G. Bhanot, W. Fishler and S. Rudaz, Nucl. Phys. B 155 (1979) 208; M. E. Peskin, Nucl. Phys. B 156 (1979) 365.

[7] T. -M. Yan, Phys. Rev. D 22 (1980) 1652.

[8] K. Shizuya, LBL preprint LBL-11004 (September, 1980).

[9] M. Shifman, A. Vainshtein and V. I. Zakharov, Nucl. Phys. B 147 (1979) 385; ibid, 448.

[10] R. Barbieri et al., Nucl. Phys. B 105 (1976) 125.

[11] M. K. Gaillard and B. W. Lee, Phys. Rev. Letters 33 (1974) 108; G. Altarelli and L. Maiani, Phys. Letters 52B (1975) 351.

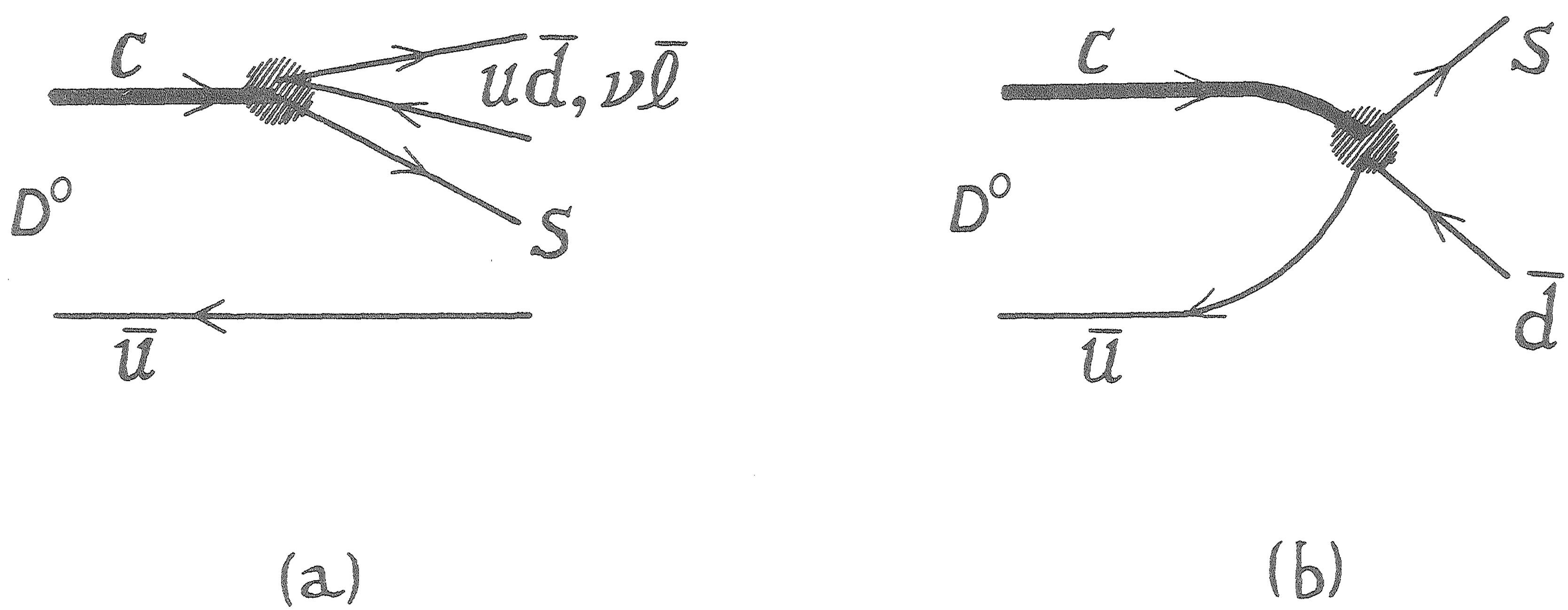

Figure 1

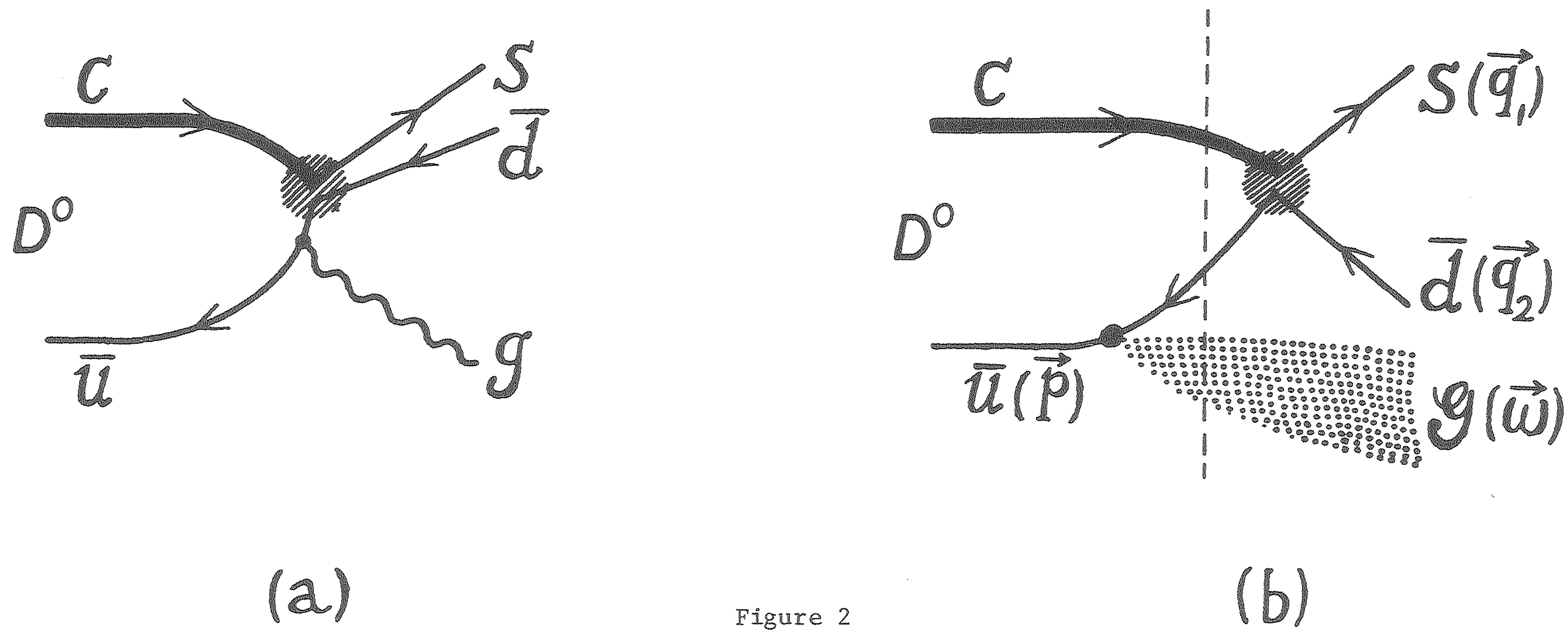

Figure 2